

The Higgs as a Portal to Plasmon-like Unparticle Excitations

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Abstract

A renormalizable coupling between the Higgs and a scalar unparticle operator \mathcal{O}_U of non-integer dimension $d_U < 2$ triggers, after electroweak symmetry breaking, an infrared divergent vacuum expectation value for \mathcal{O}_U . Such IR divergence should be tamed before any phenomenological implications of the Higgs-unparticle interplay can be drawn. In this paper we present a novel mechanism to cure that IR divergence through (scale-invariant) unparticle self-interactions, which has properties qualitatively different from the mechanism considered previously. Besides finding a mass gap in the unparticle continuum we also find an unparticle pole reminiscent of a plasmon resonance. Such unparticle features could be explored experimentally through their mixing with the Higgs boson.

1 Introduction

The very active field of unparticles grew out of two seminal papers [1] in which Georgi entertained the possibility of coupling a scale-invariant sector (with a non-trivial infrared fixed point) to our familiar standard model of particles. He described several very unconventional features of that sector that could be probed through such couplings. In his original proposal, Georgi considered only couplings through non-renormalizable operators (after integrating out some heavy messenger sector that interacts directly both with the Standard Model and the unparticle sector). Later on Shirman et al. [2] considered the possibility of coupling directly a scalar operator of unparticles \mathcal{O}_U (of scaling dimension d_U , with $1 < d_U < 2$) to the SM Higgs field through a renormalizable coupling of \mathcal{O}_U to $|H|^2$. As pointed out in [3] such coupling induces a tadpole for \mathcal{O}_U after the breaking of the electroweak symmetry and for $d_U < 2$ the value of the vacuum expectation value $\langle \mathcal{O}_U \rangle$ has an infrared (IR) divergence. That divergence should be cured before any phenomenological implications of the Higgs-unparticle coupling can be studied in a consistent way. Ref. [3] discussed a simple way of inducing an IR cutoff that would make $\langle \mathcal{O}_U \rangle$ finite. One of the main implications of such mechanism was the appearance of a mass gap, m_g , of electroweak size for the unparticle sector. Needless to say, such mass gap has dramatic implications both for phenomenology and for constraints on the unparticle sector.

In addition, Ref. [3] showed that, after electroweak symmetry breaking (EWSB), the Higgs state mixes with the unparticle continuum in a way reminiscent of the Fano-Anderson model [4], familiar in solid-state and atomic physics as a description of the mixing between a localized state and a quasi-continuum. When the Higgs mass is below m_g , the Higgs survives as an isolated state but with some unparticle admixture that will modify its properties. On the other hand, the unparticle continuum above m_g gets a Higgs contamination that can make it more accessible experimentally. When the Higgs mass is above m_g the Higgs state gets subsumed into the unparticle continuum with the Higgs width greatly enlarged by the unparticle mixing. Such behaviour is similar to that found when the Higgs mixes with a quasi-continuum of graviscalars [5]. In both cases, with m_h above or below m_g , the properties of the mixed Higgs-unparticle system can be described quite neatly through a spectral function analysis.

The organization of the paper is as follows: after describing the previous IR problem (Section 2) we present an alternative stabilization mechanism for $\langle \mathcal{O}_U \rangle$ (Section 3). This mechanism has significant differences with respect to that used in Ref. [3]: although it also induces an unparticle mass gap¹ it involves a scale-invariant self-coupling of unparticles only and leads to the appearance of a peculiar resonance in the unparticle continuum that is reminiscent of a plasmon excitation (Section 4). The mixing between the unparticle states and the Higgs boson after EWSB gives a handle on the structure of the unparticle continuum. This is best seen in terms of an spectral function analysis which we develop in Section 5. We present our conclusions in Section 6.

¹One expects such mass gap as a generic feature of any mechanism that solves the IR problem.

2 The Infrared Problem

We start with the following scalar potential

$$V_0 = m^2 |H|^2 + \lambda |H|^4 + \kappa_U |H|^2 \mathcal{O}_U , \quad (2.1)$$

where the first two terms are the usual SM Higgs potential and the last term is the Higgs-unparticle coupling, with κ_U having mass dimension $2 - d_U$. As usual, the quartic coupling λ would be related in the SM to the Higgs mass at tree level by $m_{h^0}^2 = 2\lambda v^2$. We write the Higgs real direction as $Re(H^0) = (h^0 + v)/\sqrt{2}$, with $v = 246$ GeV.

The unparticle operator \mathcal{O}_U has dimension d_U , spin zero and its propagator is [1, 6]

$$P_U(p^2) = \frac{A_{d_U}}{2 \sin(\pi d_U)} \frac{i}{(-p^2 - i\epsilon)^{2-d_U}}, \quad A_{d_U} \equiv \frac{16\pi^{5/2}}{(2\pi)^{2d_U}} \frac{\Gamma(d_U + 1/2)}{\Gamma(d_U - 1)\Gamma(2d_U)} . \quad (2.2)$$

When the Higgs field gets a non zero vacuum expectation value (VEV) the scale invariance of the unparticle sector is broken [2]. From (2.1) we see that in such non-zero Higgs background the physical Higgs field mixes with the unparticle operator \mathcal{O}_U and also a tadpole appears for \mathcal{O}_U itself which will therefore develop a non-zero VEV.

As done in Ref. [3], it is very convenient to use a deconstructed version of the unparticle sector as proposed in [7]. One considers an infinite tower of scalars φ_n ($n = 1, \dots, \infty$) with masses squared $M_n^2 = \Delta^2 n$. The mass parameter Δ is small and eventually taken to zero, limit in which one recovers a (scale-invariant) continuous mass spectrum. As explained in [7], the deconstructed form of the operator \mathcal{O}_U is

$$\mathcal{O} \equiv \sum_n F_n \varphi_n , \quad (2.3)$$

where F_n is chosen as

$$F_n^2 = \frac{A_{d_U}}{2\pi} \Delta^2 (M_n^2)^{d_U-2} , \quad (2.4)$$

so that the two-point correlator of \mathcal{O} matches that of \mathcal{O}_U in the $\Delta \rightarrow 0$ limit. In the deconstructed theory the unparticle scalar potential, including the coupling (2.1) to the Higgs field, reads

$$\delta V = \frac{1}{2} \sum_n M_n^2 \varphi_n^2 + \kappa_U |H|^2 \sum_n F_n \varphi_n . \quad (2.5)$$

A non-zero VEV, $\langle |H|^2 \rangle = v^2/2$, triggers a VEV for the fields φ_n :

$$v_n \equiv \langle \varphi_n \rangle = -\frac{\kappa_U v^2}{2M_n^2} F_n , \quad (2.6)$$

thus implying, in the continuum limit,

$$\langle \mathcal{O}_U \rangle = -\frac{\kappa_U v^2}{2} \int_0^\infty \frac{F^2(M^2)}{M^2} dM^2 , \quad (2.7)$$

where

$$F^2(M^2) = \frac{A_{d_U}}{2\pi} (M^2)^{d_U-2} , \quad (2.8)$$

is the continuum version of (2.4). We see that $\langle \mathcal{O}_U \rangle$ has an IR divergence for $d_U < 2$, due to the fact that for $M \rightarrow 0$ the tadpole diverges while the mass itself, that should stabilize the unparticle VEV, goes to zero.

In Ref. [3] it was shown how one can easily get an IR regulator in (2.8) by including a coupling

$$\delta V = \zeta |H|^2 \sum_n \varphi_n^2 , \quad (2.9)$$

in the deconstructed theory. This coupling respects the conformal symmetry but will break it when H takes a VEV. Now one gets

$$\langle \mathcal{O}_U \rangle = -\frac{\kappa_U v^2}{2} \int_0^\infty \frac{F^2(M^2)}{M^2 + \zeta v^2} dM^2 , \quad (2.10)$$

which is obviously finite for $1 < d_U < 2$ and reads explicitly

$$\langle \mathcal{O}_U \rangle = -\frac{1}{2} \kappa_U \frac{A_{d_U}}{2\pi} \zeta^{d_U-2} v^{2d_U-2} \Gamma(d_U-1) \Gamma(2-d_U) . \quad (2.11)$$

Implications for EWSB of such coupling (2.9) were studied in Ref. [3].

3 An Alternative Solution to the IR Problem

It is natural to attempt to solve the IR problem of the previous section by introducing a quartic coupling term for the deconstructed scalar fields φ_n so that the VEVs v_n are under control. As pointed out already in Ref. [3] the naive try with $\delta V = \lambda_U \sum_n \varphi_n^4$ fails. Here we prove that the particular combination

$$\delta V = \frac{1}{4} \xi \left(\sum_{n=1}^\infty \varphi_n^2 \right)^2 , \quad (3.1)$$

is successful in providing a finite value for $\langle \mathcal{O}_U \rangle$. Before showing that explicitly, let us first show that the coupling (3.1) has a finite and scale-invariant continuum limit.

We can take as scale transformations for the deconstructed fields φ_n

$$\varphi_n(x) \rightarrow a \varphi_n(xa) , \quad (3.2)$$

while leaving the space-time coordinates unscaled ($x \rightarrow x$). It is straightforward to show that under such scale transformation the kinetic part of the (deconstructed) action is invariant while the mass terms are not, as usual. In the continuum limit, however, taking $\Delta \cdot u(M^2, x)$ as the continuum limit of $\varphi_n(x)$, and using the scale transformation

$$u(M^2, x) \rightarrow u(M^2/a^2, xa) , \quad (3.3)$$

the continuum action

$$S = \int d^4x \int_0^\infty dM^2 \left[\frac{1}{2} \partial_\mu u(M^2, x) \partial^\mu u(M^2, x) - M^2 u^2(M^2, x) \right] , \quad (3.4)$$

is indeed scale invariant. Using the same construction, it is then straightforward to see that the continuum limit of the quartic coupling (3.1) is well defined and scale invariant, being explicitly given by:

$$\delta S = - \int d^4x \int_0^\infty dM_1^2 \int_0^\infty dM_2^2 \frac{1}{4} \xi u^2(M_1^2, x) u^2(M_2^2, x) . \quad (3.5)$$

To keep the following analysis general, we consider both couplings ζ and ξ simultaneously, writing for the deconstructed part of the scalar potential:

$$\delta V = \frac{1}{2} \sum_n M_n^2 \varphi_n^2 + \kappa_U |H|^2 \sum_n F_n \varphi_n + \zeta |H|^2 \sum_n \varphi_n^2 + \frac{1}{4} \xi \left(\sum_{n=1}^\infty \varphi_n^2 \right)^2 . \quad (3.6)$$

The minimization equation for the Higgs field is not affected by the new coupling ξ , while that for $v_n \equiv \langle \varphi_n \rangle$ can be put in the form

$$v_n = \frac{-\frac{1}{2} \kappa_U v^2 F_n}{M_n^2 + \zeta v^2 + \xi \sum_{m=1}^\infty v_m^2} . \quad (3.7)$$

Squaring the above equation and summing in n from 1 to ∞ one gets an implicit equation for

$$\sigma^2 \equiv \sum_{n=1}^\infty v_n^2 . \quad (3.8)$$

In the continuum limit, and using

$$(\mu_U^2)^{2-d_U} \equiv \kappa_U^2 \frac{A_{d_U}}{2\pi} , \quad (3.9)$$

the equation for σ^2 reads

$$\sigma^2 = \frac{1}{4} (\mu_U^2)^{2-d_U} v^4 \int_0^\infty dM^2 \frac{(M^2)^{d_U-2}}{(M^2 + \zeta v^2 + \xi \sigma^2)^2} . \quad (3.10)$$

or, performing the integral explicitly,

$$\sigma^2 = \frac{1}{4} \Gamma(d_U - 1) \Gamma(3 - d_U) (\mu_U^2)^{2-d_U} v^4 (\zeta v^2 + \xi \sigma^2)^{d_U-3} , \quad (3.11)$$

which can be solved for σ^2 (numerically if $\zeta \neq 0$ or analytically if $\zeta = 0$).

The induced mass gap in the unparticle continuum is now

$$m_g^2 = \zeta v^2 + \xi \sigma^2 , \quad (3.12)$$

and it is clear that this mass gap will cutoff the IR divergence of \mathcal{O}_U even for $\zeta = 0$, solving therefore the infrared problem. Note that $\sigma \neq 0$ only if $v \neq 0$ so that the mass gap is in any case associated with EWSB.

4 Unparticle Plasmon Excitation

We begin by writing down explicitly the infinite mass matrix that mixes the (real) neutral component h^0 of the Higgs with the deconstructed tower of unparticle scalars, φ_n . The different matrix elements are:

$$M_{hh}^2 = 2\lambda v^2 \equiv m_{h0}^2, \quad (4.1)$$

$$M_{hn}^2 = \kappa_U v F_n \frac{M_n^2 + \xi \sigma^2}{M_n^2 + m_g^2} \equiv A_n, \quad (4.2)$$

$$\begin{aligned} M_{nm}^2 &= (M_n^2 + m_g^2)\delta_{nm} + \frac{1}{2}\kappa_U^2 \xi v^4 \frac{F_n F_m}{(M_n^2 + m_g^2)(M_m^2 + m_g^2)} \\ &\equiv (M_n^2 + m_g^2)\delta_{nm} + a_n a_m. \end{aligned} \quad (4.3)$$

It is a simple matter to obtain the hh -entry of the inverse (infinite matrix) propagator associated to this infinite mass matrix. Already taking its continuum limit we obtain:

$$iP_{hh}(p^2)^{-1} = p^2 - m_{h0}^2 + J_2(p^2) - \frac{1}{2}\xi v^2 \frac{[J_1(p^2)]^2}{1 + \frac{1}{2}\xi v^2 J_0(p^2)}, \quad (4.4)$$

where we have used the integrals

$$\begin{aligned} J_k(p^2) &\equiv \int_0^\infty G_U(M^2, p^2) (M^2 + \xi \sigma^2)^k dM^2 \\ &= \frac{v^2}{p^4} \left(\frac{\mu_U^2}{m_g^2} \right)^{2-d_U} \Gamma(d_U - 1) \Gamma(2 - d_U) \left\{ \left(1 - \frac{p^2}{m_g^2} \right)^{d_U-2} (p^2 - m_g^2 + \xi \sigma^2)^k \right. \\ &\quad \left. - \left[1 + (2 - d_U) \frac{p^2}{m_g^2} \right] (\xi \sigma^2 - m_g^2)^k - k p^2 (\xi \sigma^2 - m_g^2)^{k-1} \right\}, \end{aligned} \quad (4.5)$$

with integer k and where $G_U(M^2, p^2)$ is:

$$G_U(M^2, p^2) \equiv \frac{v^2 (\mu_U^2/M^2)^{2-d_U}}{(M^2 + m_g^2 - p^2)(M^2 + m_g^2)^2}. \quad (4.6)$$

These integrals are real for $p^2 < m_g^2$ but they develop an imaginary part for $p^2 > m_g^2$. This imaginary part will be important later on when we discuss the spectral function associated to $P_{hh}(p^2)$. The final expression for the inverse propagator with all the integrals explicitly performed is lengthy and not very illuminating. Although the integrals in (4.4) diverge for $p^2 \rightarrow m_g^2$, the combination entering (4.4) is finite.

In contrast with the scenario analyzed in Ref. [3], in which the (real part of the) Higgs-unparticle propagator had a pole associated with a Higgs (with non-standard couplings), the propagator (4.4) has an additional pole associated with the unparticle continuum. In order to understand the origin of this additional pole consider the unparticle submatrix

(4.3). It has a simple form (a diagonal part plus a rank-1 correction) that allows one to find a particularly interesting eigenvalue ω_{p0}^2 (and eigenvector $\{r_n\}$) that satisfy

$$1 + \sum_n \frac{a_n^2}{M_n^2 + m_g^2 - \omega_{p0}^2} = 0 , \quad (4.7)$$

and

$$r_n = \frac{a_n}{N_p(\omega_{p0}^2 - M_n^2 - m_g^2)} , \quad (4.8)$$

where N_p is a normalization constant that ensures $\sum_{n=1}^{\infty} r_n^2 = 1$. For sufficiently large values of the a_n 's Eq. (4.7) has a solution, with $\omega_{p0}^2 > m_g^2$ necessarily. Note that this pole can exist due to the presence of the new quartic coupling ξ and only after EWSB, which gives $a_n \neq 0$. The appearance of this state out of the unparticle continuum is reminiscent of the appearance of plasmon excitations in condensed matter physics. In fact, the structure of the unparticle submatrix is similar to the Hamiltonian that describes the residual long-range Coulomb interactions induced in a plasma by a probe electromagnetic wave. Such structure lies at the root of different collective phenomena in different fields of physics [8].

The previous discussion can be carried over to the continuum limit, in which the condition (4.7) takes the form

$$1 + \frac{1}{2}\xi v^2 \text{P.V.} \left[\int_0^\infty G_U(M^2, \omega_{p0}^2) dM^2 \right] = 1 + \frac{1}{2}\xi v^2 \text{P.V.}[J_0(\omega_{p0}^2)] = 0 , \quad (4.9)$$

and will be modified only quantitatively by the mixing of unparticles with the Higgs in the full matrix (4.1)-(4.3). In general we will expect two poles, one Higgs-like at m_h^2 coming from the unmixed m_{h0}^2 , and one plasmon-like at ω_p^2 coming from the unmixed ω_{p0}^2 , both of them somewhat displaced by the mixing.

5 Spectral Function Analysis

In order to study in more detail this interplay between the Higgs and the unparticle sector it is instructive to examine the spectral representation of the mixed propagator (4.4), which is given by

$$\rho_{hh}(s) = -\frac{1}{\pi} \text{Im}[P_{hh}(s + i\epsilon)] , \quad (5.1)$$

where the limit $\epsilon \rightarrow 0$ is understood. We can easily calculate this spectral function by using $1/(x + i\epsilon) \rightarrow \text{P.V.}[1/x] - i\pi\delta(x)$ directly in the integrals J_k of (4.5) to obtain, for $s > m_g^2$,

$$J_k(s + i\epsilon) = R_k(s) + iI_k(s) , \quad (5.2)$$

with

$$\begin{aligned}
R_k(s) &= \frac{v^2}{s^2} \left(\frac{\mu_U^2}{m_g^2} \right)^{2-d_U} \Gamma(d_U - 1) \Gamma(2 - d_U) \left\{ \left(\frac{s}{m_g^2} - 1 \right)^{d_U-2} (s - m_g^2 + \xi \sigma^2)^k \cos(d_U \pi) \right. \\
&\quad \left. - \left[1 + (2 - d_U) \frac{s}{m_g^2} \right] (\xi \sigma^2 - m_g^2)^k - k s (\xi \sigma^2 - m_g^2)^{k-1} \right\} , \\
I_k(s) &= \pi \frac{v^2}{s^2} (s + \xi \sigma^2 - m_g^2)^k \left(\frac{\mu_U^2}{s - m_g^2} \right)^{2-d_U} .
\end{aligned} \tag{5.3}$$

As in the case of Ref. [3] there are two qualitatively different cases, depending on whether the Higgs mass m_h is larger or smaller than m_g . For $m_h < m_g$, the spectral function is explicitly given by

$$\rho_{hh}(s) = \frac{1}{K^2(m_h^2)} \delta(s - m_h^2) + \theta(s - m_g^2) \frac{T_U(s)}{\mathcal{D}^2(s) + \pi^2 T_U^2(s)} , \tag{5.4}$$

where $\mathcal{D}(s)$ and $T_U(s)$ are the real and imaginary parts of $iP_{hh}(s + i\epsilon)^{-1}$ when $s > m_g^2$:

$$iP_{hh}(s + i\epsilon)^{-1} = \mathcal{D}(s) + i T_U(s) . \tag{5.5}$$

More explicitly, one finds

$$\mathcal{D}(s) = s - m_{h0}^2 + R_2(s) \tag{5.6}$$

$$- \frac{1}{2N(s)} \xi v^2 \left\{ \left[1 + \frac{1}{2} \xi v^2 R_0(s) \right] \left[R_1(s)^2 - I_1(s)^2 \right] + \xi v^2 I_0(s) R_1(s) I_1(s) \right\} ,$$

$$T_U(s) = \frac{v^2}{s^2 N(s)} (s - \xi \sigma^2 - m_g^2)^2 \left(\frac{\mu_U^2}{s - m_g^2} \right)^{2-d_U} , \tag{5.7}$$

with

$$N(s) \equiv \left[1 + \frac{1}{2} \xi v^2 R_0(s) \right]^2 + \left[\frac{1}{2} \xi v^2 I_0(s) \right]^2 . \tag{5.8}$$

Finally

$$K^2(s_0) \equiv \frac{d}{ds} \mathcal{D}(s) \Big|_{s=s_0} . \tag{5.9}$$

An explicit expression for $K^2(s_0)$ can be obtained directly from $\mathcal{D}(s)$ above, but we do not reproduce it here.

One can check that the spectral function (5.4) is properly normalized:

$$\int_0^\infty \rho_{hh}(s) ds = 1 . \tag{5.10}$$

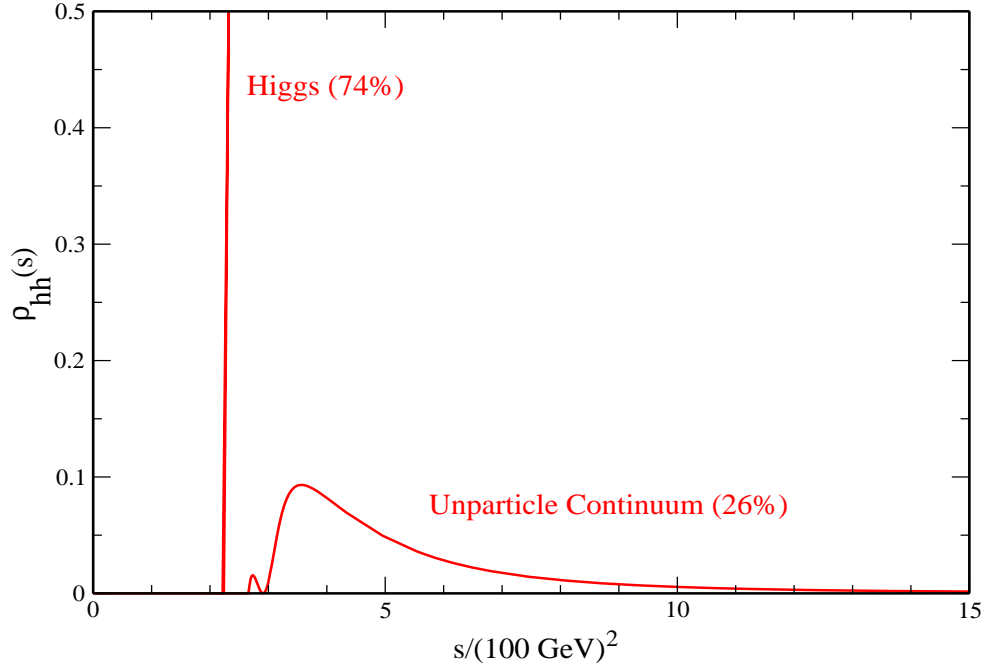


Figure 1: Spectral function with a Higgs below m_g , obtained for the case $\zeta = 0.4$, $\xi = 0.1$, $m^2 = 0$, $\kappa_U = v^{2-d_U}$ and $d_U = 1.2$. The percentage of Higgs composition of the isolated pole and of the unparticle continuum is given in parenthesis.

The physical interpretation of this spectral function is the standard one: Let us call $|h\rangle$ the Higgs interaction eigenstate and $|u, M\rangle$ the unparticle interaction eigenstates (a continuous function of M) and $|H\rangle$, $|U, M\rangle$ the respective mass eigenstates after EWSB. Then one has

$$\begin{aligned}
 |\langle H|h\rangle|^2 &= \frac{1}{K^2(m_h^2)} , \\
 |\langle U, M|h\rangle|^2 &= \frac{T_U(M^2)}{\mathcal{D}^2(M^2) + \pi^2 T_U^2(M^2)} ,
 \end{aligned} \tag{5.11}$$

so that ρ_{hh} describes in fact the Higgs composition of the isolated pole and the unparticle continuum. The proper normalization (5.10) is simply a consequence of the proper normalization of $|h\rangle$, *i.e.* $|\langle h|h\rangle|^2 = 1$. From the simple form of $T_U(s)$ in (5.7) we can see directly that for $M_0^2 = m_g^2 + \xi\sigma^2$ the spectral function is zero, corresponding to an unparticle state $|U, M_0\rangle$ which has $\langle h|U, M_0\rangle = 0$. The amount of $|h\rangle$ admixture in any state is important because it will determine key properties of that state, like its coupling to gauge bosons, that are crucial for its production and decay.

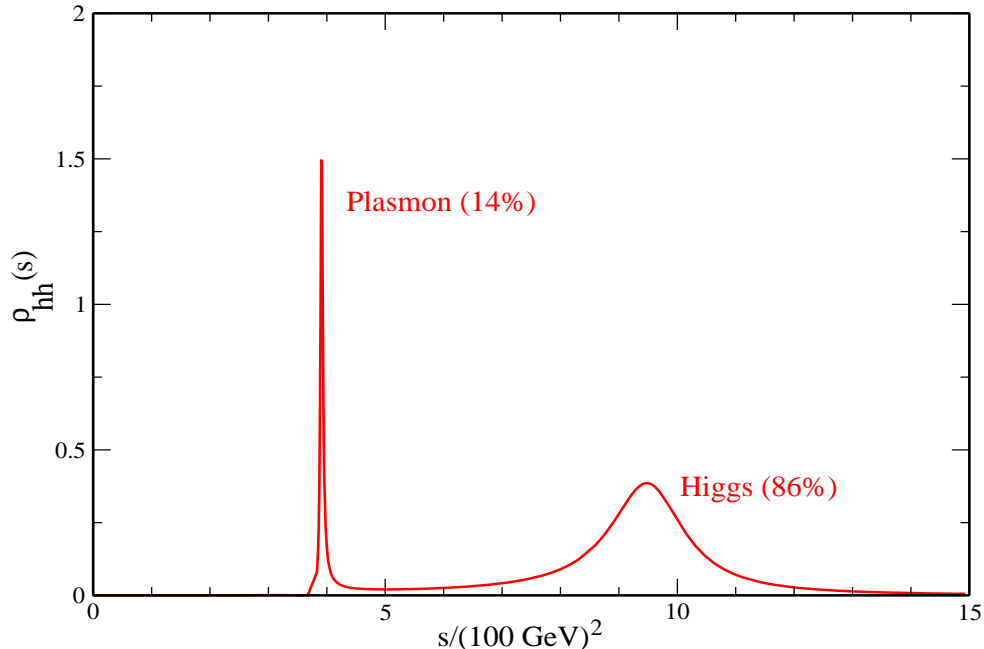


Figure 2: Spectral function with plasmon and Higgs above m_g , obtained for the case $\zeta = 0.3$, $\xi = 0.2$, $m^2 = -1.5(100 \text{ GeV})^2$, $\kappa_U = v^{2-d_U}$ and $d_U = 1.1$. The percentage of Higgs composition of each resonance is given in parenthesis.

Fig. 1 shows the spectral function for a case with $m_h < m_g$. The parameters have been chosen as follows: $d_U = 1.2$, $\kappa_U = v^{2-d_U}$, $m^2 = 0$, $\zeta = 0.4$ and $\xi = 0.1$. We see a Dirac delta at $m_h^2 = (152 \text{ GeV})^2$, a mass gap for the unparticle continuum at $m_g^2 = (163 \text{ GeV})^2$, and a zero at $M_0^2 = (171 \text{ GeV})^2$. There is also a plasmon-like resonance at $\omega_p^2 = (176 \text{ GeV})^2$ but it is not very conspicuous in this particular case. In parenthesis we give the percentage of Higgs composition in the isolated resonance and in the continuum: it is simply given by the integral of $\rho_{hh}(s)$ in the corresponding region. We see that the Higgs has lost some of its original Higgs composition due to mixing with the unparticles (as in the usual singlet dilution) while the unparticle continuum gets the lost Higgs composition spread above m_g (in a continuum way reminiscent of the models considered in [9]).

The plasmon-like resonance can be seen much more clearly in other cases, like the one shown in Fig. 2, which has $m_h > m_g$. It corresponds to $d_U = 1.1$, $\kappa_U = v^{2-d_U}$, $m^2 = -1.5(100 \text{ GeV})^2$, $\zeta = 0.3$ and $\xi = 0.2$ and has a mass gap at $m_g^2 = (164 \text{ GeV})^2$, a Higgs resonance at $m_h^2 = (307 \text{ GeV})^2$ and a plasmon-like spike at $\omega_p^2 = (198 \text{ GeV})^2$. There is also a zero at $M_0^2 = (188 \text{ GeV})^2$ right below the plasmon resonance, but it cannot be

discerned in the plot due to the scale of the figure. We give again in parenthesis the Higgs composition of the Higgs and plasmon resonances. For $m_h > m_g$, the spectral function is given by the second part of (5.4) only, without a Dirac delta-function, and there is no separate $|H\rangle$ state.

The shape of the continuum around the resonances at $s_r = \{m_h^2, \omega_p^2\}$ can be obtained directly from the spectral density (5.4) by writing

$$\mathcal{D}(s) \simeq (s - s_r) K^2(s_r) , \quad (5.12)$$

where $K^2(s_r)$ is defined in Eq. (5.9). In this case, with $s_r > m_g^2$, one should be careful about using the principal value definition of the integrals entering $\mathcal{D}(s)$ to properly calculate its derivative at s_r . Substituting (5.12) in the spectral function (5.4), we see that the resonances have a Breit-Wigner shape of width Γ_r given by

$$\frac{\Gamma_r}{\sqrt{s_r}} = \frac{\pi T_U(s_r)}{s_r K^2(s_r)} . \quad (5.13)$$

6 Conclusions

An unparticle sector could be explored experimentally in a very interesting way if it is coupled to the Standard Model directly through the Higgs $|H|^2$ operator. In this paper we have revisited such couplings of the Higgs to an unparticle scalar operator \mathcal{O}_U of non-integer dimension d_U . We have expanded upon our previous work [3] by considering a new way of solving the infrared problem that affects the expectation value of \mathcal{O}_U for $d_U < 2$ [3] that is generated by EWSB. We have shown how a scale-invariant unparticle self-coupling² can in fact generate a mass gap m_g for unparticles that acts as an IR cutoff to give a finite $\langle \mathcal{O}_U \rangle$.

In addition to solving the IR problem, the new coupling can induce after EWSB a new resonance in the unparticle continuum through a mechanism quite similar to those giving rising to plasmon resonances in condensed matter systems [8]. The mass mixing of unparticles with the Higgs after EWSB results in a spectrum of states with some admixture of Higgs that will dictate some of their production and decay properties. One can distinguish two generic types of spectra. In the first, there is an isolated state below the mass gap, which one would typically identify with the Higgs boson although it will carry some unparticle admixture that will change its properties with respect to a SM Higgs (e.g. the coupling to gauge bosons will be reduced). Beyond the mass gap there will be an unparticle continuum (possibly with a large plasmon resonance) that will be accessible experimentally through its Higgs admixture.

In the second type of spectrum, the Higgs mass will be above the mass gap and the Higgs resonance will in fact merge with the unparticle continuum acquiring a significant width. In addition to this resonance a large plasmon resonance can also be present. Both

²The importance of unparticle self-interactions for phenomenology has been emphasized in [10].

resonances will have some Higgs admixture so that both could show up experimentally as Higgses with non-standard properties.

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